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Recovering Relativistic Nuclear Phenomenology from the Quark-Meson Coupling Model

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Abstract

The quark-meson coupling (QMC) model for nuclear matter, which describes nuclear matter as non-overlapping MIT bags bound by the self-consistent exchange of scalar and vector mesons is modified by the introduction of a density dependent bag constant. The density dependence of the bag constant is related to that of the in-medium effective nucleon mass through a scaling ansatz suggested by partial chiral symmetry restoration in nuclear matter. This modification overcomes drawbacks of the QMC model and leads to the recovery of the essential features of relativistic nuclear phenomenology. This suggests that the modification of the bag constant in the nuclear medium may play an important role in low- and medium-energy nuclear physics.

Since quantum chromodynamics (QCD) is believed to be the correct theory underlying strong interactions, the physics of nuclei is, in essence, an exercise in applied QCD. Building connections between observed nuclear phenomena and the interactions and symmetries of the underlying quark and gluon degrees of freedom has become one of the principal goals of nuclear theorists. However, due to the complexity of the low-energy QCD, knowledge of QCD has had very little impact, to date, on the study of low- and medium-energy nuclear phenomena.

A reasonable consensus is that the relevant degrees of freedom for nuclear physics at low energy scales are hadrons instead of quarks and gluons. One general approach, relativistic nuclear phenomenology which has gained much credibility during last twenty years, is to treat nucleons in nuclear environment as point-like Dirac particles interacting with large canceling scalar and vector potentials. This approach has been successful in describing the spin-observables of nucleon-nucleus scattering in the context of relativistic optical potentials [1,2]. Moreover, such potentials can be derived from the relativistic impulse approximation [2]. The relativistic field-theoretical models based on nucleons and mesons, QHD, also feature Dirac nucleons interacting through the exchange of scalar and vector mesons [3]. QHD, at the mean-field level, has proven to be a powerful tool for describing the bulk properties of nuclear matter and spin-orbit splittings of finite nuclei [3]. It is known that the large and canceling scalar and vector potentials are central to the success of the relativistic nuclear phenomenology. Recent progress in understanding the origin of these large potentials for propagating nucleons in nuclear matter has been made via the analysis of the finite-density QCD sum rules [4].

Given the wide success of the Dirac approach in describing various low- and medium-energy nuclear phenomena and the support from the finite-density QCD sum rules, it is a challenge to study the relevance of the quark structure of the nucleon to the dynamics of normal nuclei. A few years ago, Guichon [5] proposed a quark-meson coupling (QMC) model to investigate the direct "quark effects" in nuclei. In this model, nuclear matter consists of

non-overlapping MIT bags interacting through the self-consistent exchange of mesons in the mean-field approximation, and the mesons are directly coupled to the quarks. This simple QMC model has been refined by including nucleon Fermi motion and the center-of-mass corrections to the bag energy [6] applied to variety of problems [7–13]. (There have been several works that also discuss the quark effects in nuclei, based on other effective models for the nucleon [14]).

Although it provides a simple and attractive framework to incorporate the quark structure of the nucleon in the study of nuclear phenomena, the QMC model has a serious shortcoming. It predicts much smaller scalar and vector potentials for the nucleon than obtained in relativistic nuclear phenomenology and finite-density QCD sum rules. Unless there is a large isoscalar anomalous coupling (ruled out by other considerations) this implies a much smaller nucleon spin-orbit force in finite nuclei. To lowest order in the nucleon velocity and potential depth the nucleon spin-orbit potential can be obtained in a model independent way from the strengths of the scalar and vector potentials. The spin-orbit potential from the QMC model is too weak to successfully explain spin-orbit splittings in finite nuclei and the spin-observables in nucleon-nucleus scattering.

We observe that the bag constant is held to be at its free space value in the QMC model for nuclear matter. This assumption can be questioned. In the MIT bag model, the bag constant denotes the vacuum energy (relative to the perturbative vacuum), which contributes ~ 300 MeV to the nucleon energy and provides the necessary pressure to confine the quarks. Thus, the bag constant is an inseparable ingredient of the bag picture of a nucleon. When a nucleon bag is put into the nuclear medium, the bag as a whole reacts to the environment. As a result, the bag constant may be modified.

There is little doubt that at sufficiently high densities, the bag constant is eventually melted away and quarks and gluons become the appropriate degrees of freedom. It, thus, seems reasonable that the bag constant be modified and decrease as density increases. Moreover, it is argued in Ref. [15] that the MIT bag constant is related to the energy associated

with the chiral symmetry restoration (the vacuum energy difference between the chiral-symmetry-restored vacuum inside the bag and the broken phase outside). Since chiral symmetry is partially restored in nuclear medium [16], the in-medium bag constant should drop relative to its free space value [15,16]. This physics has been ignored in the QMC model.

In this Letter, we shall introduce the in-medium modification of the bag constant in the QMC model for nuclear matter. We relate the in-medium bag constant to the in-medium effective nucleon mass through a scaling ansatz suggested by partial chiral symmetry restoration in nuclear matter. We find that the essential features of the relativistic nuclear phenomenology, in particular the large canceling scalar and vector potentials and hence strong spin-orbit force for the nucleon, can be recovered when the decrease of the bag constant in medium is taken into account. This suggests that the drop of the bag constant in nuclear medium relative to its free space value may play an important role in low- and medium-energy nuclear physics.

In the QMC model, the nucleon in nuclear medium is assumed to be a static spherical MIT bag in which quarks interact with the scalar and vector fields, $\overline{\sigma}$ and $\overline{\omega}$, and these fields are treated as classical fields in the mean field approximation. (Here we only consider up and down quarks.) The quark field, $\psi_q(x)$, inside the bag then satisfies the equation of motion:

$$\left[i\,\partial\!\!\!/ - (m_q^0 - g_\sigma^q\,\overline{\sigma}) - g_\omega^q\,\overline{\omega}\,\gamma^0\right]\,\psi_q(x) = 0\ , \tag{1}$$

where m_q^0 is the current quark mass, and g_σ^q and g_ω^q denote the quark-meson coupling constants. The normalized ground state for a quark in the bag is given by [5–7]

$$\psi_q(t,\mathbf{r}) = \mathcal{N}\,e^{-i\epsilon_q t/R}\begin{pmatrix} j_0(xr/R) \\ i\,\beta_q\,\sigma\cdot\hat{\mathbf{r}}\,j_1(xr/R)\end{pmatrix}\frac{\chi_q}{\sqrt{4\pi}}\ , \tag{2}$$

where $\epsilon_q = \Omega_q + g_\omega^q\,\overline{\omega}\,R$ and $\beta_q = \sqrt{(\Omega_q - R\,m_q^*)/(\Omega_q + R\,m_q^*)}$, with $\Omega_q \equiv \sqrt{x^2 + (R\,m_q^*)^2}$, $m_q^* = m_q^0 - g_\sigma^q\,\overline{\sigma}$, R the bag radius, and χ_q the quark spinor. The normalization factor is

given by $\mathcal{N}^{-2} = 2\,R^3\,j_0^2(x)\left[\Omega_q(\Omega_q-1)+R\,m_q^*/2\right]/x^2$. The x value is determined by the boundary condition at the bag surface, $j_0(x) = \beta_q\,j_1(x)$.

The energy of a static bag consisting of three ground state quarks can be expressed as

$$E_{\rm bag} = 3\,\frac{\Omega_q}{R} - \frac{Z}{R} + \frac{4}{3}\,\pi\,R^3\,B\ , \tag{3}$$

where Z is a parameter which accounts for zero-point motion and B is the bag constant. In the calculations to follow, we use R_0, B_0 and Z_0 to denote the corresponding bag parameters for the free nucleon. After the corrections of spurious c.m. motion in the bag, the effective mass of a nucleon bag at rest is taken to be [6,7]

$$M_N^* = \sqrt{E_{\rm bag}^2 - \langle p_{\rm cm}^2\rangle}\ , \tag{4}$$

where $\langle p_{\rm cm}^2\rangle = \sum_q \langle p_q^2\rangle$ and $\langle p_q^2\rangle$ is the expectation value of the quark momentum squared, $(x/R)^2$.

The QMC model assumes that both Z and B are independent of density. (The bag radius is determined by the equilibrium condition for the bag, see below.) This assumption is unjustified. As argued in Ref. [15], the MIT bag constant is essentially the energy associated with the chiral symmetry restoration. Therefore, one expects the bag constant to drop relative to its free space value as a consequence of partial chiral symmetry restoration in the nuclear medium. Of course, one has to invoke model descriptions in order to obtain a quantitative estimate for the reduction of the bag constant. According to the scaling ansatz advocated by Brown and Rho [17], the bag constant should scale like $B/B_0 \simeq \Phi^4$ [15,16]. Here Φ denotes the universal scaling, and $\Phi \sim m_\rho^*/m_\rho \simeq f_\pi^*/f_\pi \simeq (M_N^*/M_N)^{2/3}$ has been suggested in Ref. [16]. Motivated by these suggestions, we introduce the following scaling ansatz

$$\frac{B}{B_0} = \left(\frac{M_N^*}{M_N}\right)^{\kappa/3}\ , \tag{5}$$

for the in-medium bag constant. The case $\kappa = 0$ corresponds to the usual QMC model. Note that combining Eqs. (3), (4), and (5) yields a self-consistency condition for B. In

principle, the parameter Z may also be modified in the nuclear medium. However, unlike the bag constant, it is unclear how Z changes with the density as Z is not directly related to chiral symmetry. Here we assume that the medium modification of Z is small at low and moderate densities and take $Z = Z_0$.

The bag radius is determined by the equilibrium condition for the bag, $\partial M_N^*/\partial R = 0$. In free space, one may fix M_N at its experimental value 939 MeV and use the equilibrium condition to determine the bag parameters. For several choice of bag radius, $R_0 = 0.6, 0.8, 1.0$ fm, the results for $B_0^{1/4}$ and Z_0 are $188.1, 157.5, 136.3$ MeV and $2.030, 1.628, 1.153$, respectively.

The total energy per nucleon at finite density, ρ_N, can be written as [7]

$$E_{\rm tot} = \frac{\gamma}{(2\pi)^3\,\rho_N}\int^{k_F} d^3k\sqrt{M_N^{*2}+\mathbf{k}^2} + \frac{g_\omega^2}{2m_\omega^2}\,\rho_N + \frac{m_\sigma^2}{2\,\rho_N}\overline{\sigma}^2 \ , \tag{6}$$

where γ is the spin-isospin degeneracy, and $\gamma = 4$ for symmetric nuclear matter and $\gamma = 2$ for neutron matter. Here we have used that the mean field $\overline{\omega}$ created by uniformly distributed nucleons is determined by baryon number conservation to be [5–7]

$$\overline{\omega} = \frac{3\,g_\omega^q\,\rho_N}{m_\omega^2} = \frac{g_\omega\,\rho_N}{m_\omega^2} \ , \tag{7}$$

where $g_\omega \equiv 3g_\omega^q$. The scalar mean field is determined by the thermodynamic condition $(\partial\, E_{\rm tot}/\partial\,\overline{\sigma})_{R,\rho_N} = 0$, which yields the self-consistency condition

$$g_\sigma\,\overline{\sigma} = \frac{g_\sigma^2}{m_\sigma^2}\,C(\overline{\sigma})\frac{\gamma}{(2\pi)^3}\int^{k_F} d^3k\,\frac{M_N^*}{\sqrt{M_N^{*2}+\mathbf{k}^2}} \ , \tag{8}$$

where $g_\sigma \equiv 3\,g_\sigma^q$ and

$$C(\overline{\sigma}) = \frac{E_{\rm bag}}{M_N^*}\left[\left(1-\frac{\Omega_q}{E_{\rm bag}\,R}\right)\,S(\overline{\sigma}) + \frac{m_q^*}{E_{\rm bag}}\right]\left[1-\frac{\kappa}{3}\,\frac{E_{\rm bag}}{M_N^{*2}}\frac{4}{3}\,\pi\,R^3\,B\right]^{-1} \ , \tag{9}$$

with

$$S(\overline{\sigma}) = \frac{\Omega_q/2 + R\,m_q^*\,(\Omega_q-1)}{\Omega_q\,(\Omega_q-1) + R\,m_q^*/2} \ . \tag{10}$$

We now turn to present numerical results. For simplicity, we take $m_q^0 = 0$. The coupling constants g_σ and g_ω are chosen to fit the nuclear matter binding energy (-16 MeV) at

TABLE I. Coupling constants and nuclear incompressibility K (in MeV) at various κ values. The case of $\kappa = 0$ corresponds to the simple QMC model and the last raw gives the result of QHD-I. Here we have used $m_\sigma = 550$ MeV and $m_\omega = 783$ MeV.

κ value	R_0=0.6 fm			R_0=0.8 fm			R_0=1.0 fm		
	$g_\sigma^2/4\pi$	$g_\omega^2/4\pi$	K	$g_\sigma^2/4\pi$	$g_\omega^2/4\pi$	K	$g_\sigma^2/4\pi$	$g_\omega^2/4\pi$	K
$\kappa = 0$	20.17	1.56	223	21.87	1.14	200	22.47	0.96	189
$\kappa = 7$	4.40	4.41	356	4.51	3.51	314	4.54	3.18	304
$\kappa = 8$	3.13	5.72	426	3.14	4.50	359	3.14	4.07	350
$\kappa = 9$	2.10	8.32	650	2.04	6.33	471	2.00	5.67	415
$\kappa = 10$	1.21	14.89	2233	1.20	11.17	1058	1.16	9.74	918
QHD-I	8.45	12.84	540	8.45	12.84	540	8.45	12.84	540

the saturation density ($\rho_{\rm N}^0$ =0.17 fm^{-3}). The resulting coupling constants and the nuclear incompressibility are listed in Table 1. We note that while the scalar coupling decreases, the vector coupling and the nuclear incompressibility increase as κ increases.

This result can be understood from the scaling ansatz Eq. (5). When $\kappa = 0$, the strength of the vector field is much smaller than that required in QHD-I. This, in Refs. [6,7], is attributed to the repulsion provided by the c.m. corrections to the bag energy. When $\kappa > 0$, Eq. (5) provides a new source of attraction as it effectively reduces M_N^*. Consequently, additional vector field strength is required to balance the nuclear matter. The decrease of the scalar coupling with increasing κ is due to the increasingly strong attraction from the dropping bag constant. (When $\kappa > 10$, the self-consistent solution around $\rho_N = \rho_N^0$ does not exist as the attraction gets too strong.) The rapid increase of the nuclear incompressibility with increasing κ results from that the contribution of the vector field to the nuclear incompressibility is proportional to g_ω^2.

The total energy per nucleon for symmetric nuclear matter is presented in Fig. 1 for various κ values, with $R_0 = 0.8$ fm. The result from QHD-I is also plotted for comparison.

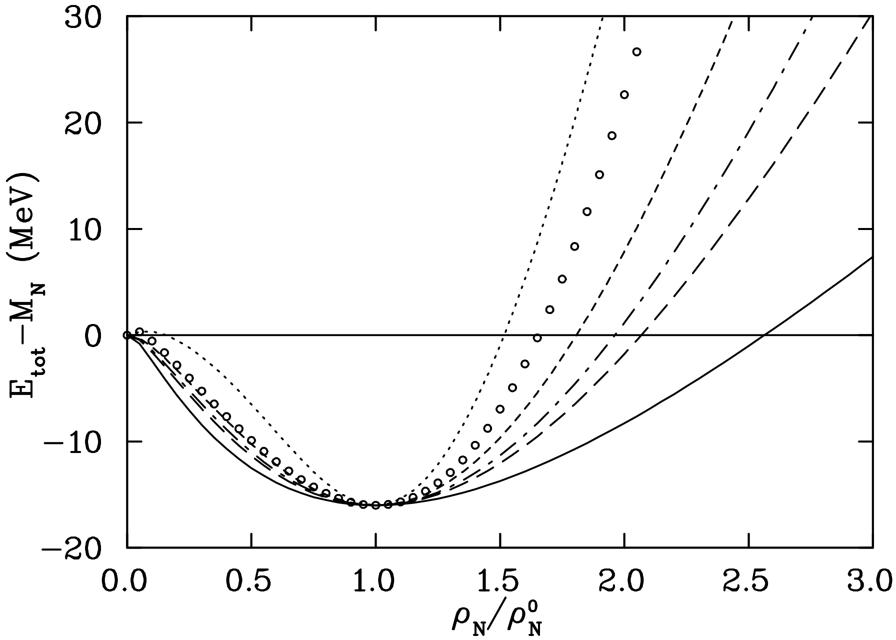

FIG. 1. Energy per nucleon for symmetric nuclear matter as a function of the medium density, with $R_0 = 0.8$ fm. The five curves correspond to $\kappa = 0$ (solid), 7 (long-dashed), 8 (dott-dashed), 9 (short-dashed), and 10 (dotted), respectively. The result from QHD-I is given by the open circles.

The usual QMC model ($\kappa = 0$) predicts a much softer equation of state for the nuclear matter than in QHD-I. As κ gets larger, the equation of state becomes stiffer, and when $\kappa \sim 9.5$ the equation of state is essentially the same as the one predicted in QHD-I. The resulting effective mass and the vector field strength are shown in Fig. 2. One can see clearly that the effective mass decreases and the vector field strength increases rapidly as κ increases. As shown in Ref. [13], the equivalent scalar and vector potentials appearing in the wave equation for a point-like nucleon are $M_N^* - M_N$ and $U_{\rm v} \equiv g_\omega \overline{\omega}$, respectively. Thus, our results indicate that the scalar and vector potentials for the nucleon become stronger as κ gets larger.

The corresponding in-medium bag radius is shown in Fig. 3. In the QMC model, the bag radius decreases slightly in the medium. When $\kappa > 0$, the bag constant drops relative to its free space value, which implies a decreasing bag pressure and hence gives rise to a increasing bag radius in the medium. With larger κ, the bag radius grows more quickly.

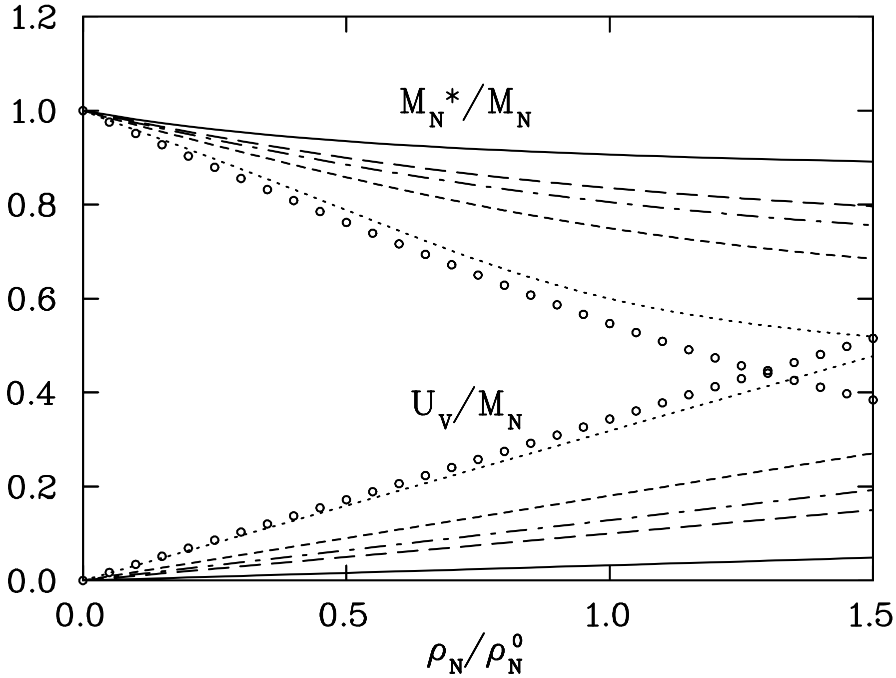

FIG. 2. Results for the ratios M_N^*/M_N and $U_\mathrm{v}/M_N \equiv g_\omega \overline{\omega}/M_N$ as functions of the medium density, with $R_0 = 0.8$ fm. The five curves correspond to, $\kappa = 0$ (solid), 7 (long-dashed), 8 (dot-dashed), 9 (short-dashed), and 10 (dotted), respectively. The result from QHD-I is given by the open circles.

This is consistent with the "swollen" nucleon picture drawn from the decrease of the meson masses in nuclear medium [17–26] (see, however, Ref. [27]). The in-medium bag constant is plotted in Fig. 4. One can see that when $\kappa > 0$ the bag constant decreases as density increases. The rate of this decrease gets larger for larger κ values. Finally, the sensitivity of our results to the free space bag radius is illustrated in Fig. 5. For a given κ value, one finds that the ratios B/B_0 and M_N^*/M_N increase and the ratio R/R_0 and the vector field strength decrease as R_0 increases.

We observe that $\kappa \sim 0$ and $\kappa \sim 7 - 10$ lead to qualitatively different physics. Unless one expresses the bag constant in terms of QCD operators and solves QCD in the nuclear matter, κ is unknown. Nevertheless, one may get an estimate from the Brown-Rho scaling ansatz, $\Phi \simeq m_\rho^*/m_\rho$. Using $m_\rho^*/m_\rho \sim 0.8$ at $\rho_N = \rho_N^0$ as suggested in Ref. [25] and

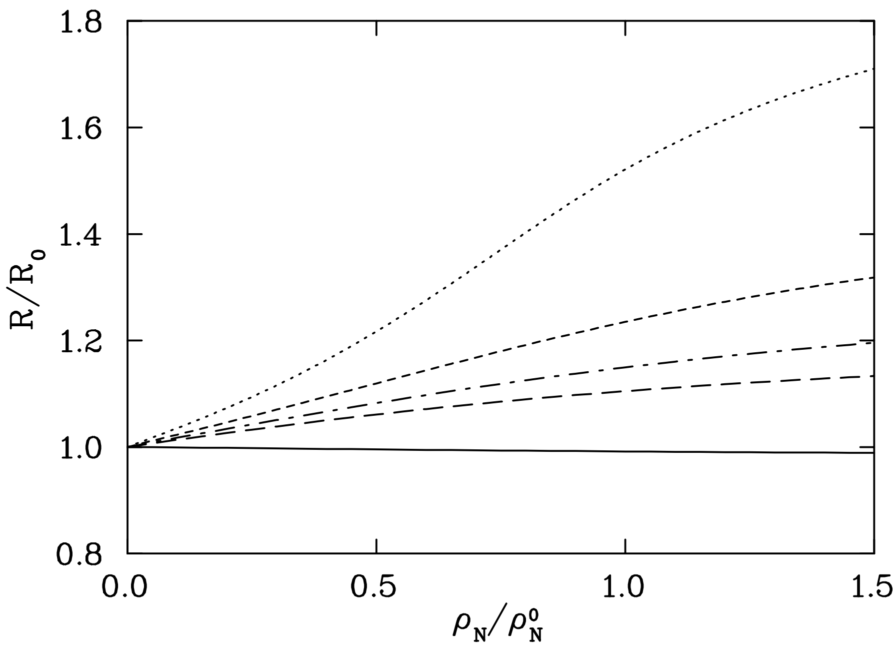

FIG. 3. Result for the ratio R/R_0 as a function of the medium density, with $R_0 = 0.8$ fm. The five curves correspond to $\kappa = 0$ (solid), 7 (long-dashed), 8 (dot-dashed), 9 (short-dashed), and 10 (dotted), respectively.

$B/B_0 \simeq \Phi^4$, we get $B/B_0 \simeq 0.4$ at the saturation density. This requires $\kappa \simeq 8.6$, 9.1, 9.3 for $R_0 = 0.6$, 0.8, 1.0 fm, respectively. For these κ values, the reduction of the bag constant dominates the attraction, the internal structure of the nucleon only plays a relatively minor role, and the large and canceling scalar and vector potentials for the nucleon appears naturally. Such potentials are comparable to those suggested by Dirac phenomenology [1,2], Brueckner calculations [2], and finite-density QCD sum rules [4], but somewhat smaller than those obtained in QHD-I. These large potentials also imply a strong nucleon spin-orbit potential. Therefore, we conclude that the essential features of relativistic nuclear phenomenology can be recovered when the decrease of the bag constant with increasing density is considered.

The QMC model is valid only if the nucleon bags do not overlap significantly. With the κ values suggested above, we find $R/R_0 \sim 1.25$ at $\rho_N = \rho_N^0$. For $R_0 = 0.6 - 0.7$ fm, as suggested by Guichon [5], one finds $R \sim 0.75 - 0.875$ fm, which gives $4\pi R^3 \rho_N^0/3 \sim 0.3 - 0.48$.

This indicates that the overlap between the bags is reasonably small at the saturation density, though a factor of two larger than in the usual QMC model. For larger R_0 and/or higher densities, the overlap becomes more significant and the non-overlapping picture of the nuclear matter may become inadequate. However, it is unclear at this stage whether the overlap between the bags is effectively included in the scalar and vector mean fields. Further study is needed to clarify this issue. We also note that for the κ values suggested above, the resulting nuclear incompressibility is comparable to that obtained in QHD-I, which is too large compared with the empirical value and that obtained in the usual QMC model. This may be fixed by introducing self-interactions of the scalar field, which, however, will introduce more free parameters.

In summary, we have included the decrease of the bag constant in the quark-meson coupling model for the nuclear matter. This effectively introduces a new source of attraction, which needs to be compensated with additional vector field strength. When the change of

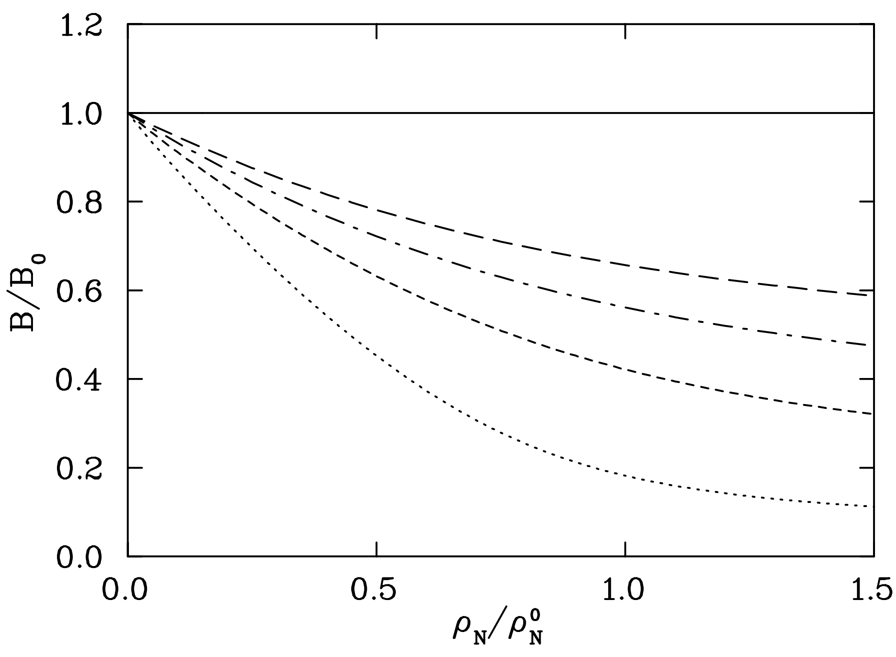

FIG. 4. Result for the ratio B/B_0 as a function of the medium density, with $R_0 = 0.8$ fm. The five curves correspond to $\kappa = 0$ (solid), 7 (long-dashed), 8 (dot-dashed), 9 (short-dashed), and 10 (dotted), respectively.

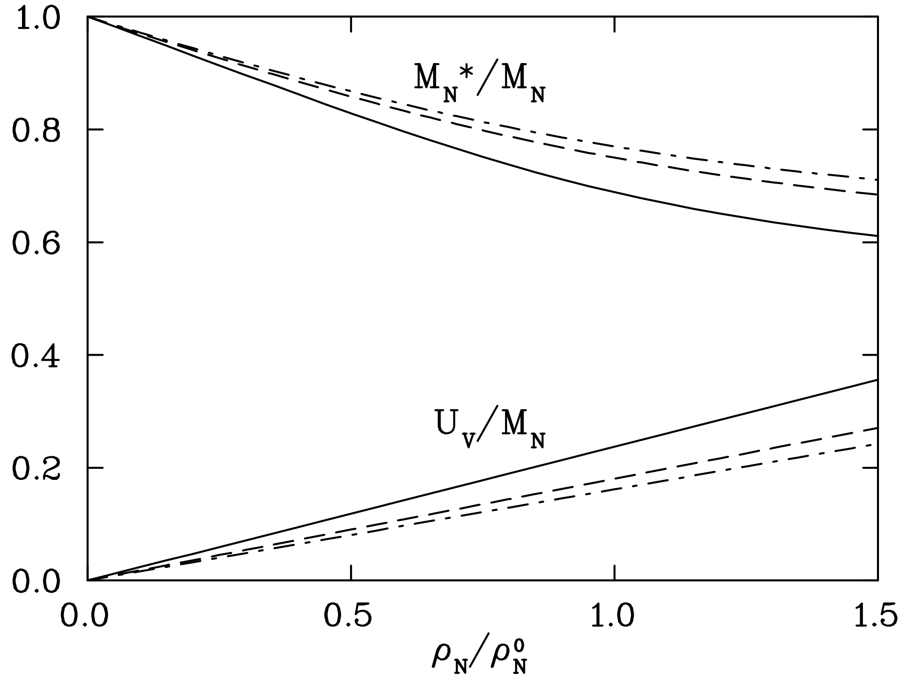

FIG. 5. Results for the ratios M_N^*/M_N and $U_{\rm v}/M_N \equiv g_\omega \overline{\omega}/M_N$ as functions of the medium density, with $\kappa = 9.0$. The three curves are for $R_0 = 0.6$ fm (solid), 0.8 fm (dashed), and 1.0 fm (dot-dashed), respectively.

the bag constant is large, as supported by partial chiral-symmetry restoration, large and canceling scalar and vector potentials for the nucleon emerge. The essential physics of the relativistic nuclear phenomenology can be recovered. The internal quark structure of the nucleon seems to play only a relatively minor role. On the other hand, the in-medium modification of the bag constant may play an important role in low- and medium-energy nuclear physics. The model presented in the present paper can be applied to variety of nuclear physics problems. Work along this direction is in progress and will be reported elsewhere.

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REFERENCES

[1] S. Hama, B. C. Clark, E. D. Cooper, H. S. Sherif, and R. L. Mercer, Phys. Rev. C 41 (1990) 2737, and references therein.

[2] S. J. Wallace, Ann. Rev. Nucl. Part. Sci. 37 (1987) 267, and references therein.

[3] B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. 16 (1986) 1, and references therein.

[4] T. D. Cohen, R. J. Furnstahl, D. K. Griegel, and X. Jin, Prog. Part. Nucl. Phys. Vol. 35 (1995) 221, and references therein; R. J. Furnstahl, X. Jin, and D. B. Leinweber, nucl-th/9511007.

[5] P. A. M. Guichon, Phys. Lett. B 200 (1988) 235.

[6] S. Fleck, W. Bentz, K. Shimizu, and K. Yazaki, Nucl. Phys. A 510 (1990) 731.

[7] K. Saito and A. W. Thomas, Phys. Lett. B 327 (1994) 9.

[8] K. Saito, A. Michels, and A. W. Thomas, Phys. Rev. C 46 (1992) R2149; K. Saito and A. W. Thomas, Nucl. Phys. A 574 (1994) 659.

[9] K. Saito and A. W. Thomas, Phys. Lett. B 335 (1994) 17.

[10] K. Saito and A. W. Thomas, Phys. Lett. B 363 (1995) 157.

[11] K. Saito and A. W. Thomas, Phys. Rev. C 51 (1995) 2757.

[12] H. Q. Song and R. K. Su, Phys. Lett. B 358 (1995) 179.

[13] P. A. M. Guichon, K. Saito, E. Rodionov, and A. W. Thomas, University of Adelaide preprint ADP-95-45/T194 (1995), nucl-th/9509034.

[14] M. K. Banerjee, Phys. Rev. C 45 (1992) 1359; V. K. Mishra, Phys. Rev. C 46 (1992) 1143; E. Naar and M. C. Birse, J. Phys. G 19 (1993) 555.

[15] C. Adami and G. E. Brown, Phys. Repts. 234 (1993) 1, and references therein.

[16] G. E. Brown, M. Buballa, Z. Li, and J. Wambach, Nucl. Phys. A 593 (1995) 295; G. E. Brown and Mannque Rho, hep-ph/9504250, and references therein.

[17] G. E. Brown and M. Rho, Phys. Rev. Lett. 66 (1991) 2720.

[18] R. Altemus *et al.*, Phys. Rev. Lett. 44 (1980) 965; P. Barreau *et al.*, Nucl. Phys. A 402 (1983) 515; Z. E. Meziani *et al.*, Phys. Rev. Lett. 52 (1984) 2130; Z. E. Meziani *et al.*, Phys. Rev. Lett. 54 (1985) 1233; A. Zghiche *et al.*, Nucl. Phys. A 572 (1994) 513.

[19] D. Reffay-Pikeroen *et al.*, Phys. Rev. Lett. 60 (1988) 776; A. Magnon *et al.*, Phys. Lett. B 222 (1989) 352; J. E. Ducret *et al.*, Nucl. Phys. A 556 (1993) 373.

[20] J. V. Noble, Phys. Rev. Lett. 46 (1981) 412.

[21] L. S. Celenza, A. Rosenthal, and C. M. Shakin, Phys. Rev. C 31 (1985) 232.

[22] G. E. Brown and M. Rho, Phys. Lett. B 222 (1989) 324.

[23] M. Soyeur, G. E. Brown, and M. Rho, Nucl. Phys. A556 (1993) 355.

[24] G. E. Brown, C. B. Dover, P. B. Siegel, and W. Wiese, Phys. Rev. Lett. 60 (1988) 2723.

[25] T. Hatsuda and S. H. Lee, Phys. Rev. C 46 (1992) R34; M. Asakawa and C. M. Ko, Nucl. Phys. A 560 (1993) 399;M. Asakawa and C. M. Ko, Phys. Rev. C 48 (1993) R526; X. Jin and D. B. Leinweber, Phys. Rev. C 52 (1995) 3344.

[26] H. Kurasawa and T. Suzuki, Nucl. Phys. A 490 (1988) 571; Prog. Theor. Phys. 84 (1990) 1030; K. Tanaka, W. Bentz, A. Arima, and F. Beck, Nucl. Phys. A 528 (1991) 676; J. C. Caillon and J. Labarsouque, Phys. Lett. B 311 (1993) 19; H. -C. Jean, J. Piekarewicz, and A. G. Williams, Phys. Rev. C 49 (1994) 1981; H. Shiomi and T. Hatsuda, Phys. Lett. B 334 (1994) 281.

[27] I. Sick, Phys. Lett. 157B (1985) 13; C. Atti, Nucl. Phys. A 479 (1989) 349c; A. Zghiche *et al.*, Nucl. Phys. A 572 (1994) 513; S. Ishikawa *et al.*, Phys. Lett. B 339 (1995) 293.